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# Scalar Black Holes

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## ABSTRACT

A class of exact regular spherically symmetric solutions to the Einstein equation obeying Dymnikova's definition  $T_1^1 = T_0^0$  of vacuumlike state is considered. These solutions, that may be interpreted as *black holes*, are not only singularity free, but also do not set us thinking about the loss of information under gravitational collapse. However, according to the singularity theorems, the geometries introduced by these solutions inevitably have some causal pathology.

If vacuumlike state is reinterpreted to be a sort of *confinement* representing a particular phase of matter, this pathology under certain conditions does not, however, involve actual causality violation. By means of that, the mentioned above class of solutions, named in what follows *scalar black holes*, and probably much broader variety of solutions including dynamic ones, may be incorporated in General Relativity. There are listed other phenomena, such as 'hidden mass', that could help to identify the presence of vacuumlike phase.

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# 1 Introduction

During the last decade, the internal structure of black holes has attracted considerable attention, with issues of mass-energy singularity and nonunitary black hole dynamics being in focus (e.g., [1, 2]). The latest development has produced the perturbative quantum models of *string* black holes with quantum microstates being countable, and closed-string Hawking emission going on in a unitary way. The emission rate agrees with the semiclassical Hawking prediction. quantum models still contain the ‘central’ singularity but in a ‘civilized’ form: though the sources of fields propagating in the black hole interior are viewed as located at singularity, the modes of the fields are well defined up to event horizon.

Despite the fact that quantum models still have some shortcomings (e.g., only integer moded states are accountable with the current perturbative technique), their very existence shows that the loss of information is not a must for black holes.

The confirmation of this remarkable discovery demands, however, the construction of a macroscopic black hole model with similar properties because a fundamental criterion of applicability of a theory still is the preservation of macroscopic causality. The violation of just this ‘principle of macroscopic censorship’ prompted J. Wheeler to name gravitational collapse ‘the greatest crises in physics of all time’ [3].

Put very simply, by their structure the standard black hole models (e.g. the Schwarzschild one) provide no capacity for storage of propagation modes and quantum numbers, which therefore are condemned to die in gravitational trap. The most efforts to eliminate this shortcoming (e.g., [4, 5]) introduce the storage capacity by proposing some quantum phenomena that collect information in the *vacuum* dynamic region of a black hole. Instead, as its starting point, the present work takes the Dymnikova model [6, 7, 8] based on the exact regular solution to the Einstein equation [8, 9] obeying the condition

$$T_1^1 = T_0^0, \quad T_2^2 = T_3^3, \quad T_k^i = 0 \quad \text{if} \quad i \neq k. \quad (1)$$

We shall consider a class of solutions obeying (1). Each of these solutions with final, but large enough, Schwarzschild mass depicts asymptotically flat spacetime with an event horizon. The presence of the horizon identifies them as black holes. Inside the event horizon, there exists a *nonsingular mass-energy distribution* that seems to be a natural carrier of the information, which supposedly disappears in Schwarzschild black hole. Thus, the imperative to admit the nonunitary development of gravitational collapse is eliminated.

The singularity theorems [10, 11, 12], which are an integral part of GR, impose, however, severe restrictions on the use of the regular black hole solutions: particles, following paths allowed by such solutions, unavoidably run into some causality violation.

The essence of difficulties is as follows. Let a nonsingular spherical *core* replace the Schwarzschild singularity. By symmetry, the metric at the core center is Minkowskian. Thus, in some vicinity of the center, a particle can move both toward the center and outward. As it turns out, a particle, freely moving outward, leaves the core, enters the dynamic region, and then traverses the outer space by crossing the event horizon. In the dynamic region,

the movement of this particle proceeds, however, against the time arrow established by the inward falling particles, which cross the horizon from the outer space. Thus, the regular black hole has an ‘oncoming time topology’ preventing consistent definition of the arrow of time in the dynamic region. According to the third Tipler theorem [12] this discordance is generic rather than conditioned by a high level of symmetry.

The widely accepted paradigm, considering geometry to be the warrant of causal evolution of universe, implies that in these circumstances one must reject either regular black hole solutions or causality (if not General Relativity itself). This paradigm, however, is not a part of GR, and looks suspicious. Indeed, the Einstein equation

$$G_{ik} = -\kappa T_{ik} \quad (2)$$

*algebraically* links two quantities, which observationally are mutually independent. Thus, the Einstein equation is an *equivalence relation* (invariant, but otherwise just like  $E = mc^2$ ). Therefore properties attributed to matter should be taken on equal terms as those related to geometry. Thus, in some circumstances just the properties of matter can protect causality whereas the geometry by itself *allows* causality violation (cf. [13]).

This modifies the view of gravitational collapse. A causal geometric deficiency may be admitted in conjunction with the states of matter, which properties prevent the actual violation of causality on their own.

Following this program, we examine, which properties of the matter involved let us incorporate the regular black hole solutions in GR. The answer amounts to three main conditions: (1) the matter is in *vacuumlike state*; (2) the vacuumlike state is a *particle confinement* of a kind; and (3) the high-density vacuumlike continuum exhibits *combined stress pattern*.

## 2 Vacuumlike continuum

Incorporating the cosmological term,  $\lambda g_{ik}$ , in his equation, Einstein [14] was the first who introduced a Lorentz-invariant, and in this sense vacuumlike, mass-energy distribution. He, however, assumed the term to be related to the background of spacetime geometry. With coming the dynamic models pioneered by Friedmann [15], Einstein threw this term out, but since then the idea of modified or modifiable vacuum surfaced repeatedly.

Along with metric tensor  $g_{ik}$ , the scalar-tensor theory of gravitation [16, 17] introduces also scalar field  $\varphi$ . Without the gravity-connected terms, the Lagrangian of scalar-tensor theory reduces to the Lagrangian for massless scalar field:

$$L = \frac{1}{2} g^{ab} \varphi_{,a} \varphi_{,b} - \varepsilon(\varphi), \quad \varepsilon(\varphi) \geq 0, \quad (3)$$

The stress-energy tensor of its homogeneous state ( $\varphi_{,i} = 0$ ,  $i = 0, 1, 2, 3$ ) is

$$T_k^i = \delta_k^i \varepsilon, \quad \varepsilon = \text{const} \geq 0. \quad (4)$$

. Possessing the structure of cosmological term, this tensor is Lorentz invariant. In this sense, the scalar field  $\varphi$  is vacuumlike. Shortly after the scalar-tensor theory appeared, it has been understood [18] that GR allows the existence of *vacuumlike continuum* that, in distinction to the cosmological term, represents a *phase of matter*. i.e., can participate in mutual phase transitions with phases of matter having stress-energy tensors of general algebraic type.

According to (4), for vacuumlike phase

$$T_0^0 = T_1^1 = T_2^2 = T_3^3, \quad (5)$$

i.e., its energy density  $\varepsilon \equiv T_0^0$  and the main pressures  $p_i \equiv -T_i^i$ ,  $i = 1, 2, 3$  (*no summation*), are related by

$$p_i = -\varepsilon, \quad i = 1, 2, 3. \quad (6)$$

Thus, the energy being positive, the pressures are negative. This seems to involve spontaneous contraction – mechanical collapse of continuum. In general relativity, however, this is not the case, and (4) corresponds to the equilibrium de Sitter solution with *negative pressures counterbalanced by the divergence of geodesics*, i.e., by the *repulsive* gravitational action. Hence, the vacuumlike phase has no classical counterpart, and is *essentially relativistic* substance comprehensible only in GR frame.

As was immediately realized, the vacuumlike phase is a candidate for the upper state of matter attainable at growing density. Sakharov [19] proposed that it is the initial cosmological state providing nonsingular onset of expanding universe. At the same time, there was offered nonsingular Friedmann cosmology [20]. Much later development brought the vacuumlike phase down to GUT and inflationary cosmologies [21, 22].

As applied to the black hole physics, a series of works was initiated by Markov [23] who proposed that quantum corrections bound the spacetime curvature by Planckian value; this value being reached, the effective stress-energy tensor takes the form (4). This hypothesis, elaborated by Frolov, Markov, and Muchanov [24], has been further studied by Poisson and Israel [25], Balbinot and Poisson [26], and Balbinot [27]. In distinction from cosmology, where (5) can be easily incorporated into existing formalisms, in case of black hole the Israel junction conditions [28, 29] require the existence of a transition layer between the core and vacuum. The layer cannot be vacuumlike in the sense of (4). Also, conceived as a thin shelf, it represents a *spacelike* surface [27]. In other words, unlike the timelike surface of ordinary body, it exists at a single instant of time [27].

This brief survey outlines the sphere, which the present work belongs to. It has been partly initiated by the essential Dymnikova's idea [6] to reject equation (5) as the exhaustive macroscopic definition of vacuumlike phase. She offers equations (1), which introduce the *minimal spherical symmetry* instead of homogeneity. She reasonably maintains that the medium obeying (1) *has no unique comoving reference frame*, so that it is a vacuumlike one. Her black hole model is based on the exact *regular* solution to the Einstein equation:

$$ds^2 = \left(1 - \frac{R_g(r)}{r}\right) dt^2 - \left(1 - \frac{R_g(r)}{r}\right)^{-1} dr^2 - r^2 d\Omega^2,$$

$$d\Omega^2 \equiv d\theta^2 + \sin^2 \theta d\varphi^2, \quad R_g(r) \equiv r_g \left( 1 - e^{-\frac{r^3}{r_0^2 r_g}} \right), \quad r_0 \equiv \sqrt{\frac{3}{8\pi\epsilon_0}} \quad (7)$$

where  $r_g$  is gravitational radius and  $\epsilon_0$  the energy density at the center of the black hole configuration. The entire family of regular solutions defined by the Dymnikova condition (1) is considered in the next Section.

Let us consider the properties of continuum obeying homogeneous conditions (5).

Generally, the stress-energy tensor,  $T_k^i$ , defines the proper mass-energy density,  $\epsilon$ , as well as the local velocity,  $u^i$ , of a physical continuum, by means of the eigen-value equation

$$T_a^i u^a = \epsilon u^i. \quad (8)$$

If the algebraic structure of  $T_k^i$  provides the vector field  $u^i$  to exist, and is timelike and unique, we will name the continuum *particlelike*. In this case the field  $u^i$  generates a unique congruence of paths, to which it is tangent. The *test particles* (or *observers*, or *fluid elements*) propagating along these paths form *comoving reference frame*. By definition, continuum is *at rest* relative to its comoving frame.

Eq. (8) may have no solution, as, e.g., in case of free electromagnetic field. If continuum obeys (5), Eq. (8) degenerates into a trivial identity  $0 = 0$ , and the notion of local velocity ceases to be applicable at all, just as it is for ordinary vacuum. This is the case of homogeneous vacuumlike continuum. Its properties, immediately evident from Eq. (5), are as follows:

1. If (5) holds, any vector  $u^i$  satisfies (8), so that the notion of the velocity of a particle relative to the continuum loses any definite meaning.
2. In analogy with ordinary vacuum, in each reference frame free falling in the continuum, the local physics, including the results of the measurements over the continuum, is the same.
3. In default of the unique comoving reference frames, the concept of Eulerian (moving with the medium) element of vacuumlike continuum cannot be established and, by means of that, also *the concept of particles that constitute this continuum*.
4. In the continuum obeying (5) everywhere, the particle creation is forbidden because of the uncertainty in the impulse of the particles that would be created. This means that the phase transition of vacuumlike continuum to a phase with particles is to be initiated by particlelike seeds.
5. Since the notion of a particle velocity relative to vacuumlike continuum cannot be established, any interactions between a particle and this continuum are independent of the particle velocity (principle of relativity).
6. This means that there is no basis for the kinetic thermal exchange between vacuumlike and particlelike continua. Thus, the concept of the temperature *in Boltzmann's sense* cannot be introduced for vacuumlike continuum, and by means of that, the concept of Boltzmann's entropy. Thus, the macroscopic internal state of vacuumlike continuum is defined by the *only parameter* – its mass-energy density.

7. Since for particlelike phase Boltzmann's temperature is well defined, the possibility of phase transition between particlelike and vacuumlike phases makes us to attribute an effective temperature to vacuumlike phase, e.g., the temperature of the particlelike phase newly created in a reversible process. In distinction from the classical temperature, this effective temperature is not *an independent variable*, but a function of vacuumlike mass-energy density (cf. Eq. (10) below and [30]).

8. Because of the divergence of geodesics inside vacuumlike continuum, two test particles, free falling on neighboring geodesics, are accelerated one relative to another. This *gravitational repulsion* is balanced by the compressing negative pressure (6), so that the continuum is in equilibrium.

A simple model of the phase transition of a (not necessary homogeneous initially) particlelike medium to a homogeneous vacuumlike state is delivered by the relativistic equation of perfect fluid motion:

$$(\varepsilon + p) \frac{\delta u^i}{\delta s} = p_{,a} (g^{ia} - u^i u^a) , \quad \varepsilon \geq 0. \quad (9)$$

Let us assume that, with growing density  $\varepsilon$ , an internal attraction (negative pressure) appears in the fluid. Then the fluid will undergo self-compression. But when the pressure  $p \rightarrow -\varepsilon$ , its gradient,  $p_{,a}$ , vanishes. Such a self-aligning process could produce a homogeneous state with pressure satisfying (6). The final state is Lorentz invariant and therefore vacuumlike; actual particles disappear as *observable* entities, and the particlelike structure is virtual.

A simple example may be useful to overview the distinction between actual and virtual behaviors. According to Gibbons and Hawking [32], an observer, falling freely in de Sitter spacetime, reveals the presence of thermal radiation of temperature

$$T = \sqrt{\frac{2}{3\pi}} \frac{\hbar c}{k_B} \sqrt{G\varepsilon} \quad (K). \quad (10)$$

When actual radiation is present, any two observers in relative motion feel it differently, whereas the value (10) is observer-independent. This means that the Gibbons-Hawking radiation is virtual, rather than actual. From this point of view, (10) indicates not the real presence of radiation but only determines the equilibrium conditions between continuum and a measuring device: in equilibrium, the device has Gibbons-Hawking temperature. As a part of the device state, the radiation captured by the device from the vacuumlike thermal bath is, of course, real. Just as we may expect, the equilibrium conditions do not depend on the velocity of the device and defined by the only parameter – the mass-energy density of continuum

Generalizing this example, we learn first that the existence of phase transitions between vacuumlike and particlelike phases is the *requirement of thermodynamics*; and, second, that this transition is conditioned by the presence of *inhomogeneity* breaking off the ban on the particle creation in homogeneous vacuumlike phase.

Thus, we come to the view of the phase transition to vacuumlike state as brought on by the rise in the particlelike phase of internal attraction becoming the dominant internal force. Particles, confined inside the final bound state, can display its presence only within the limits of the uncertainty principle, as the *virtual* particle structure.

Along with vacuum, the resulting phase resembles the phenomena of particle physics in two ways. First, it is a kind of confinement [31] similar to the quark confinement where an attraction between quarks conceals their properties as free particles. Second, the overwhelming attractive forces appear at high concentration of mass-energy. This resemblance lets us anticipate that, having much in common with the known in particle physics, the vacuumlike confinement does not violate quantum conservation laws.

A profound approach that covers the aspects mentioned above follows from the fact that the characteristic features of vacuumlike continuum are closely similar to those of ordinary vacuum except two distinctions: the non-vanishing Lorentz-invariant stress-energy tensor and the proposed ability to retain quantum charges. These distinctions are just the ones required by classical and quantum conservation laws for considering vacuum a phase of matter. If one is governed by the idea of similarity of vacuumlike phase to the ordinary vacuum, the values of quantities describing a vacuumlike phase should be considered *zero-point values* from which the corresponding values for particlelike phase should be reckoned.

This suggests that the phase transition to vacuumlike state may be considered a *shift* of the *zero-point* densities of mass- energy and quantum charges from their zero values (which are characteristic of ordinary vacuum) to non-zero levels, what provides the accommodation of the energy and quantum charges of the particles eliminated as actual physical entities. From this point of view, vacuumlike continuum is vacuum with *shifted* zero-point levels. This approach seems to be the most consistent though it is, probably, not absolutely necessary for understanding the black hole models considered.

Considering vacuumlike continuum in this way, we, supposedly, should admit that inside continuum the spectrum of particle masses can change. For our goals, it is enough to accept the intuitive expectation that the mass-energy density, naturally connected with an actual particle inside the continuum, should exceed the mass-energy density of vacuumlike phase (which we considered above to be zero-point level). Then, some light-weight particles may be nonexistent in vacuumlike continuum. We will return to this suggestion in Section 4.

Now let us turn our attention to a localized and therefore inhomogeneous distribution of vacuumlike phase. According to the Einstein equation, the externally observed mass of a localized object depends only on the internal mass-energy distribution. Therefore, its vacuumlike properties have not come to light in its external behavior as a mechanical and gravitating body. The internal equilibrium, however, essentially depends on the algebraic structure of its stress-energy tensor. One faces two main possibilities: either vacuumlike continuum exhibits only purely normal stress (as perfect fluid) or combined stress (as, e.g., elastic body). In the former case no regular steady-state inhomogeneous solutions exist. Thus, it is our guess that, at least at collapse-induced-densities, the vacuumlike phase endures *combined stress*.

Under combined stress, the stress-energy tensor ceases to be Lorentz invariant exactly. This, however, does not ruin the vacuumlike behavior rooted in the lacking of actual particles or, from a more general point of view, in the shift of the zero-point values. All the properties listed above, therefore, remain valid, at least at the scale small in comparison with the characteristic length of inhomogeneity when the mass-energy distribution must be considered as the united whole.

### 3 Scalar black holes

We shall restrict our consideration by the Dymnikova family of solutions to the Einstein equation with diagonal stress-energy tensor obeying the Dymnikova condition (1). In this case metric can be written in the form

$$ds^2 = A(r) dt^2 - \frac{dr^2}{A(r)} - r^2 d\Omega^2, \quad (11)$$

and vice versa, from (11) follows (1).

If, for a given mass-energy distribution  $\varepsilon = \varepsilon(r)$ , the Schwarzschild mass of the whole configuration

$$M = \frac{4\pi}{c^2} \int_0^\infty \varepsilon \zeta^2 d\zeta < \infty, \quad (12)$$

the solution to the Einstein equation is given by

$$A(r) = 1 - \frac{2G}{c^4} \frac{4\pi}{r} \int_0^r \varepsilon \zeta^2 d\zeta. \quad (13)$$

Thus, within the restriction (12), the solution is defined by the mass-energy distribution  $\varepsilon = \varepsilon(r)$ , which can be chosen arbitrary. The stress-energy tensor is then defined as follows:

$$T_0^0 = T_1^1 = \varepsilon, \quad T_2^2 = T_3^3 = \varepsilon + \frac{r}{2} \varepsilon_{,r}. \quad (14)$$

Note,  $T_2^2 = T_3^3 < 0$  corresponds to positive pressure. The latter may be assumed unacceptable for vacuumlike phase. According to (14), to eliminate positive pressure, the mass-energy density must drop not faster than  $r^{-2}$  everywhere, except some vicinity of  $r = 0$ .

Let us introduce dimensionless quantities:

$$\chi \equiv \frac{r}{r_g} \quad \text{and} \quad \mu \equiv \frac{\varepsilon}{\varepsilon_m}; \quad r_g \equiv \frac{2GM}{c^2}, \quad \varepsilon_m \equiv \frac{1}{4\pi} \frac{Mc^2}{r_g^3}. \quad (15)$$



For each *arbitrary chosen* smooth  $\mu = \mu(\chi)$ , the function\*

$$A(\chi) = 1 - \frac{1}{\chi} \int_0^\chi \mu \varsigma^2 d\varsigma \quad (16)$$

determines a subclass of similar spacetimes distinguished by the value of the mass  $M$ . Substituting  $k \mu(\chi)$  for  $\mu(\chi)$ , we, for each given  $\mu(\chi)$ , can find such critical value of  $k = k_{cr}$  that  $A(\chi)$  is non-negative for  $k \leq k_{cr}$ , but for  $k > k_{cr}$  changes its sign at some  $\chi = \chi_h$ . Since  $A(0) = A(\infty) = 1$ , the zeros  $\chi_h$  are emerging in pairs. Thus, there are two essentially different classes of solutions. One, with  $A(\chi)$  being everywhere non-negative, contains *starlike* solutions. Another one, distinguishing by the presence of zeros of  $A(\chi)$ , represents *black hole* solutions. As shown in the next Section, these solutions are physically acceptable only if in some vicinity of  $\chi = 0$  the matter is in vacuumlike phase; otherwise causality cannot be preserved. We will name these solutions *scalar black holes*.

The Dymnikova Eq. (7) gives an example of solutions of both classes: a starlike solution for mass-energy density  $\varepsilon_0$  being small enough, and a scalar black hole solution for larger  $\varepsilon_0$  values.

The surfaces where  $A(\chi)$  changes sign represent horizons of a kind, similar, but not identical to the Schwarzschild horizon. A scalar black hole has at least a *pair of horizons*. We will only consider black holes with one pair of horizons and distinct the *inner horizon*,  $\chi = \chi_i$ , and the *outer* one,  $\chi = \chi_e$ . On the horizons, the  $\chi, t$ -coordinates are mutually converted from spacelike to timelike or vice versa.<sup>†</sup> In the *dynamic region* between horizons, no particle can be at rest relative to  $\chi$  which is timelike there, whereas  $t$  is spacelike. In the region  $\chi_i > \chi > 0$ , in the *core*, the orientation of  $\chi, t$ -coordinates is the same as in the *outer space*  $\chi > \chi_e$ :  $\chi$  is spacelike and  $t$  timelike. By symmetry, the metric at the center  $\chi = 0$  of the core is Minkowskian; it is closed to the de Sitter one in its vicinity. Since  $\mu(\chi)$  is an arbitrary function, the configuration of scalar black holes can have a variety of patterns with horizons that can be located inside or outside the sphere  $r = r_g$ .

## 4 Challenge to causality

The structure of a scalar black hole, dominated by the presence of pairs of horizons, is radically distinct from the Schwarzschild black hole. Two phenomena, the *oncoming time topology* and *quasistatic mass-energy distribution* in the dynamic region, challenge the very understanding of causality in GR.

To a certain extent, scalar black hole recalls the Schwarzschild one. A ‘*testing observer*’, radially falling onto this black hole from the outer space, reaches the outer horizon in a finite proper time and enters black hole interior. From this moment, in analogy with the

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\*The use of notation ‘ $A(\chi)$ ’ that is not correct in strong mathematical sense should not lead to misunderstanding.

<sup>†</sup>This fact, of course, does not interfere with using these coordinates for analytic description.

Schwarzschild case, he perceives his journey into diminishing  $r$  as changes in the environment in the course of time - until, again in a final proper time, he has found himself on the inner horizon (instead of singularity). Recalling the argumentation addressed to the singularity in Schwarzschild black hole, one might erroneously conclude that, like the Schwarzschild singularity, the core lies in the absolute future of the outer horizon. With observer's deepening into the core, the Schwarzschild-like scenario is, however, over.

Free falling into or out of a scalar black hole (11) is determined by equations:

$$r_{,t}^2 = c^2 A^2 \left(1 - \frac{A}{b}\right) ; \quad r_{,tt} = c^2 A A_{,r} \left(1 - \frac{3}{2} \frac{A}{b}\right), \quad (17)$$

$$r_{,s}^2 = b - A, \quad b = const ; \quad r_{,ss} = -\frac{1}{2} A_{,r} \cdot CR \quad (18)$$

Eqs. (18) define the velocity,  $r_{,t}$ , and the acceleration,  $r_{,tt}$ , of an observer. Since  $A = 0$  on the horizons, a test body approaching the inner horizon from the core becomes invisible from the inside. A similar effect is detected by an outer observer keeping his eye on a body falling on the outer horizon (well known black hole feature). As shown below, these horizons cannot be, nevertheless considered as one-way membranes.

Velocity,  $r_{,s}$ , and acceleration,  $r_{,ss}$ , in units of the proper time of a free falling observer are given by (18). The general pattern is as follows. If the integral of motion  $b < 1$ , the equation  $A(r) - b = 0$  has two zeros  $r = \vec{r}$  and  $r = \tilde{r}$ . One, say  $\vec{r}$ , lies in the outer space, another,  $r$ , *inside the core*. At  $\vec{r}$  and  $\tilde{r}$  observer's velocity is vanishing. At  $\vec{r}$  he is accelerated toward the black hole (quasi-Newtonian attraction), but inside the core, where  $A_{,r} < 0$ , outward. Thus, from  $r = \vec{r}$  he begins his fall into black hole, slows down inside the core, stops at  $r = \tilde{r}$ , accelerates outward, and returns to  $r = \vec{r}$  for a final proper time. Then, he repeats this cycle. If  $b > 1$ , the observer flies through the black hole and escapes to infinity. If  $b = 1$ , he gets an unstable rest at the core center; any disturbance pushes him back toward outer space. The problem is: crossing the dynamic region on his way to the core, the observer travels in the direction of time, which is opposite to the one on his way back.

Thus, on geometric basis, the arrow of time in the dynamic region between horizons cannot be chosen uniquely, this is a *NOT- region* (not orientable in time). Using the term coined apparently by Wald [11], scalar black hole has *oncoming time topology*. By means of that its geometry cannot secure the course of physical events from causality violation.

Such geometries customary are considered unacceptable. However, Riemannian geometry is alien to the notion of direction to be a distinctive geometric feature. The geometry as a bearer or a guard of the arrow of time cannot therefore be a primary notion and is only justified as a synonymic to some non-geometric circumstances (cf. Sec. 1).

The arrow of time is usually viewed to be connected with irreversible cause-and-effect sequence of events in particlelike phase. This view cannot be translated into vacuumlike phase because a sequence of causally connected events represents a unique reference frame that cannot come to existence inside vacuumlike phase. Deeper insight on the reason for this provides the uncertainty principle, which restricts the life-span of virtual particles just

enough to make it insufficient to pass or acquire definite information and by means of that let virtual particles be converted to actual ones.

The inability of vacuumlike phase to distribute information clearly manifests itself in the absence of sound. According to (6), the velocity of sound in vacuumlike phase  $c_s = c (d\varepsilon/dp)^{1/2}$  has an imaginary value.

Thus, for the core to influence the arrow of time in the dynamic region, it must contain an internal carrier of the time arrow, i.e., particlelike phase with developed cause-and-effect relations between particles. If no particles exist inside the core, would they be particles created there or penetrated from the outside, the source of oncoming time topology is lacking.

In the absence of particles serving as a reference system, the spontaneous particle creation in vacuumlike phase is forbidden. Another way for a particle to appear in the core is its penetration from the dynamic region. Since the interaction between particle and vacuumlike phase does not depend on particle's velocity, a particle inside vacuumlike phase is in the state of free falling and can leave the core. But as we discussed it in Sec. 2, a changeover from ordinary vacuum to vacuumlike phase can change the particle identity, primarily its mass. Say, only the excess of effective density of particle over the vacuumlike phase density contributes to observed particle's mass, so that in the limit of equal densities a particle is captured by the phase.

Our guess is: for a particle to exist inside a vacuumlike phase, the mass density associated with particle (say  $\rho_{pl} \sim c^3 m_{pl}^4 / \hbar^3$  with  $m_{pl}$  being the particle mass) must exceed that of vacuumlike phase. Then, vacuumlike phase is free of particles if its mass-energy density is large enough and the spectrum of particles is bounded from above (what is now commonly accepted). If the top quark with mass  $\sim 150 m_p$  is considered to be the heaviest particle, a vacuumlike continuum is free of particlelike phase if its mass density is larger than

$$c^3 \hbar^{-3} (150 m_p)^4 \sim 10^{26} g cm^{-3}. \quad (19)$$

This value may be overestimated because quark does not exist as a free particle.

Interestingly enough, nothing makes us bring Planck density in. This implies that the scalar black hole theory is in no need for quantum gravity; GR and the tools of the standard particle physics seem to be sufficient for the further development.

Thus, if the density inside the core is not less than some critical density, e.g. (19), there is no emission from the core, and the arrow of time in the dynamic region is established by matter falling from the outside. Actual bodies *fall down* into such black hole, rather than '*fall up*' (in contrast to the fictitious 'testing observer').

This way of establishing the arrow of time is seemingly indeed realized in the scalar black holes originating in gravitational collapse if the bulk transition to vacuumlike phase takes place at the critical or a larger density. Then the formation of vacuumlike core begins not earlier than the critical density is attained, and, proceeding layer by layer, keeps the density of each layer being not less than the critical one. Having been formed, the core is completely filled with vacuumlike phase, because its boundary – the inner horizon – is

immersed in vacuumlike phase. Indeed, according to (13), the distribution  $\varepsilon = \varepsilon(r)$  must continue beyond the inner horizon into the dynamic region, otherwise the horizons, and the black hole itself, do not appear. The core, therefore, does not contain cavities where particles falling from the outside could be turned back to the dynamic region.

Thus, we expect that breaching causality by particles moving outward are *geometrically admissible*, but not allowed *by the conditions on matter involved*, so that the effect of on-coming topology is nullified [33].

Another challenging feature of scalar black holes appears just because the mass-energy distribution  $\varepsilon = \varepsilon(r)$  continues beyond the inner horizon into the dynamic region.

Reasoning naively, one may speak about an element of the mass-energy distributed in the dynamic region, and, applying (18), find that each such element falls on the horizon in a final proper time. Should this be the case, the dynamic region would be free of mass-energy in a final time of any local observer.

In application to the vacuumlike phase, this, however, is wrong. The easiest way to show that is to recall that the notion of the local velocity is inapplicable to this phase, so that the movement originating in any given initial velocity distribution is unrelated to this phase. This implies that, the phase cannot be dragged in the motion along paths, which the particlelike phase follows. Thus, as soon as a vacuumlike mass-energy distribution obeys the Einstein equation, we should acknowledge the possibility of its presence in the dynamic region.

We may take a look at all that from the different point of view. Since the concept of vacuumlike phase generalizes the notion of ordinary vacuum, we should regard the mass-energy density of vacuumlike phase as the local *zero-point energy density* (Sec. 2). This returns us to the idea of a particle as primarily an *excess* in density above this zero-point level, with this excess to be the only carrier of properties of actual particles. The excess, i.e. particlelike phase, is washed out from the dynamic region, whereas the zero-point energy density distribution is, strictly speaking, just the object, which we have named ‘scalar black hole’.

Thus, our consideration of the macroscopic aspects of the stationary scalar black hole models suggests that a black hole, which is free of singularity, causality violation, and the loss of quantum information, is theoretically possible.

An essential shortcoming of our approach is the restriction by the stationary case. This is justified by mathematical simplicity rather than by physical reasoning. A realistic model of spherically symmetric scalar black hole may happen to incorporate the interaction of black hole with its vicinity not only by means of gravity and Hawking emission, but also due to the vacuumlike phase propagation. The features of the Dymnikova family of solutions seemingly supports the last possibility because vacuumlike phase in these solutions is present far from the spot where the phase transition under collapse is expected to take place.

## 5 In search of vacuumlike phase

Since the black hole interior is unavailable for observation, let us make some suggestions on searching for vacuumlike phase in other spots.

The phase transition to vacuumlike phase is the result a catastrophic self-compression, which enormously changes the magnitude of pressures. This leads us to believe that in terms of the field theory vacuumlike phase corresponds to a deep local minimum of effective potential. Then the barrier penetration to particlelike phase may be slow enough for the existence of long-living metastable vacuumlike formations that may produce observable phenomena. Even the state of the ordinary vacuum can be unattainable (just as the absolute zero of temperature). At the ‘bottom of the universe’ we will then find the low density vacuumlike phase. This phase is not required to be homogeneous as ordinary vacuum, homogeneous by definition, so that it is reasonable to speak about vacuumlike *clouds*. They may influence the observation in several ways.

First, the clouds act gravitationally. For an *outside* observer, a cloud represents hidden mass. At the same time, the cosmological objects *inside* a cloud undergo outward accelerations. These effects can misshape the pattern of homogeneous cosmological accelerations.

Second, as a whole, vacuumlike clouds behave as ordinary bodies, so that they can be in orbital movements round gravitating bodies or fall on them forming a halo. The latter must disturb the orbital movement of bodies revolving round the central body.

Third, a vacuumlike cloud may serve as ‘anti-gravity’ (diverging) lens. Distorting the background of faint galaxies, this can create the illusion of void.

Fourth, a vacuumlike cloud can be dense enough for the phase transition to particlelike state to create particle-antiparticle pairs. The subsequent particle-antiparticle annihilation will then produce characteristic annihilation electromagnetic emission from this cloud.

Fifth, the presence of vacuumlike phase can influence the output of nuclear reactions by means of capturing light particles, e.g. neutrino.

Sixth, the assumed ability of vacuumlike phase to hold quantum charges can induce some subtle differences between particles and antiparticles, e.g., such as between  $K^\circ$  and  $\bar{K}^\circ$ .

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